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Millisecond proto-magnetars and gamma ray bursts(*)

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Summary. — In the seconds after core collapse and explosion, a thermal driven-neutrino wind emerges from the cooling, deleptonizing newly-born neutron star. If the neutron star has a large-scale magnetar-strength surface magnetic field and millisecond rotation period, then the wind is driven primarily by magneto-centrifugal slinging, and only secondarily by neutrino interactions. The strong magnetic field forces the wind to corotate with the stellar surface and the neutron star's rotational energy is efficiently extracted. As the neutron star cools, and the wind becomes increasingly magnetically-dominated, the outflow becomes relativistic. Here I review the millisecond magnetar model for long-duration gamma ray bursts and explore some of the basic physics of neutrino-magnetocentrifugal winds. I further speculate on some issues of collimation and geometry in the millisecond magnetar model.

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1. – Introduction

A successful core-collapse supernova (SN) leaves behind a hot deleptonizing proto-neutron star (PNS) that cools and contracts on its Kelvin-Helmholtz cooling timescale ($\tau_{\text{KH}} \sim 10$ s), radiating its gravitational binding energy ($\sim 10^{53}$ ergs) in neutrinos [1]. A small fraction of these neutrinos deposit their energy in the tenuous atmosphere of the PNS through the interactions $\nu_e n \rightarrow pe^-$, $\bar{\nu}_e p \rightarrow ne^+$, $\nu\bar{\nu} \rightarrow e^+e^-$, and $\nu(e^-e^+) \rightarrow \nu'(e^-e^+)$. Inverse processes provide cooling. Net neutrino heating drives a thermal wind that emerges into the post-supernova shock environment [2]. For typical non-rotating non-magnetic (NRNM) neutron stars, the total kinetic energy of the wind over τ_{KH} is of order $\sim 10^{48}$ ergs and, hence, the addition to the asymptotic SN energetics is small on the scale of the canonical SN energy, $E_{\text{SN}} \sim 10^{51}$ ergs [3-5].

Magnetars —a class of neutron stars with surface magnetic field strengths of $B_0 \sim 10^{15}$ G— are thought to be born with millisecond rotation periods, their intense fields having been generated by an efficient dynamo [6-8]. Millisecond rotation periods imply a reservoir of rotational energy that is large on the scale of E_{SN} : $E_{\text{Rot}} \sim 2 \times$

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10^{52} ergs $M_{1.4} R_{\nu_{10}}^2 P_1^{-2}$, where $M_{1.4} = M/1.4 M_\odot$, $R_{\nu_{10}} = R_\nu/10$ km, and P_1 is the spin period in units of 1 millisecond (ms). Stellar progenitors that produce millisecond magnetars (MSMs) must have iron cores rotating with periods of $\lesssim 10$ s at the moment of collapse [9]. The character of collapse, bounce, and explosion can be modified by rotation. For example, rotational support leads to characteristically lower neutrino luminosity (L_ν) and longer τ_{KH} [10, 11]. In addition, a fraction of the gravitational binding of rotating collapse may be stored in shear energy and tapped by viscous processes; for parameters appropriate to MSM birth this effect can yield explosions in models that would otherwise fail [11]. Although rotation may be important during collapse and explosion, and although small-scale magnetic fields may be important in providing a viscosity capable of tapping the shear energy generated during collapse, even the magnetic energy density associated with a dipole field strength of 10^{15} G is small with respect to the thermal pressure exterior to any PNS in the first ~ 1 – 2 seconds after collapse. Therefore, similar to NRNM PNSs, we expect that the wind that accompanies MSM cooling is driven by neutrino heating at early times. However, as L_ν decreases and the thermal pressure exterior to the MSM decreases, the region exterior to the MSM must become magnetically-dominated.

The strong magnetic field forces the matter composing the outflow into near corotation with the stellar surface out to $\sim R_A$, the Alfvén point, where the magnetic energy density equals the kinetic energy density of the outflow. For $P \sim 1$ ms, if $R_A \gtrsim 15$ km, then the wind is driven primarily by magneto-centrifugal slinging; neutrino heating becomes relatively unimportant in determining the asymptotic wind velocity. Rotational energy is transferred directly from the MSM to the wind and this provides an efficient mechanism for spindown [12, 13]. The spin period of the MSM e -folds on the spindown timescale $\tau_J \approx (2/5)(M/\dot{M})(R_\nu/R_A)^2$, where \dot{M} is the mass loss rate. Because $E_{\text{Rot}} \gg E_{\text{SN}}$, just one e -folding of Ω is sufficient to modify the dynamics of the SN remnant significantly. If τ_J is small with respect to the time for the SN shock to traverse the progenitor (\sim tens of seconds for type-Ib, -Ic progenitors) we also expect this extra energy source to modify the SN nucleosynthesis [14]. As the MSM cools and the outflow becomes increasingly magnetically-dominated, R_A increases. It cannot do so indefinitely. R_A approaches the radius of the light cylinder $R_L = c/\Omega \simeq 48P_1$ km asymptotically. As it does so, the flow becomes increasingly relativistic. This is the transition between non-relativistic mass-loaded outflow and relativistic Poynting-flux dominated neutron star wind. All neutron stars, regardless of their initial spin period and magnetic field strength, go through this transition. MSMs are interesting because this transition occurs at high wind kinetic luminosity. Because E_{Rot} is large and the spindown timescale is short, and because the velocity of the wind must eventually become relativistic, these objects are a natural candidate for the central engines of long-duration gamma ray bursts (GRBs) [15, 16, 14].

2. – Proto-magnetar spindown

In this section I summarize the results of ref. [14]. See that work for more details.

Angular momentum conservation implies that $\dot{J} = d/dt(I\Omega) = -\dot{M}\mathcal{L}$, where \mathcal{L} is the specific angular momentum carried by the wind and I is the moment of inertia. In the classic model for solar spindown constructed by ref. [17], the wind problem is treated in one spatial dimension and in the equatorial plane. Consideration of the azimuthal momentum equation together with Faraday’s law gives $\mathcal{L} = R_A^2 \Omega$. To estimate the angular momentum loss rate and, thereby, the wind luminosity and the spindown timescale, we must first estimate R_A . The location of the Alfvén point depends on the radial dependence of the poloidal magnetic field. Because models with purely monopole

fields generally over-estimate spindown and models with pure dipole fields under-estimate spindown, we parameterize $B_r = B_0(R_\nu/r)^\eta$, where $2 \lesssim \eta \lesssim 3$. Taking $B^2/8\pi \sim \rho v_r^2/2$ at R_A , $\rho = \dot{M}/4\pi r^2 v_r$, and assuming that $v_A = v_r(R_A) \sim v_\phi(R_A) \sim R_A \Omega$, we find that

$$(1) \quad R_A = B_0^{2/(2\eta-1)} R_\nu^{2\eta/(2\eta-1)} (\dot{M}\Omega)^{-1/(2\eta-1)},$$

where v_A is the radial Alfvén speed, v_r is the radial velocity, v_ϕ is the azimuthal velocity, ρ is the mass density, and B_r is the radial magnetic field. Equation (1) assumes that $R_A \Omega \gg v_\nu$, where v_ν is the asymptotic wind velocity in a NRNM outflow ($v_\nu \lesssim 3 \times 10^9 \text{ cm s}^{-1}$). The absolute value of the rotational energy loss rate can be written as

$$(2) \quad \dot{E}_{\text{NR}} = B_0^{4/(2\eta-1)} R_\nu^{4\eta/(2\eta-1)} \dot{M}^{(2\eta-3)/(2\eta-1)} \Omega^{(4\eta-4)/(2\eta-1)}.$$

For parameters appropriate to MSMs

$$(3) \quad \dot{E}_{\text{NR}}^{\eta=2} \simeq 1.5 \times 10^{51} B_{0.15}^{4/3} R_{\nu_{10}}^{8/3} \dot{M}_{-3}^{1/3} P_1^{-4/3} \text{ ergs s}^{-1}$$

$$(4) \quad \dot{E}_{\text{NR}}^{\eta=3} \simeq 4.5 \times 10^{50} B_{0.15}^{4/5} R_{\nu_{10}}^{12/5} \dot{M}_{-3}^{3/5} P_1^{-8/5} \text{ ergs s}^{-1}.$$

The subscript ‘NR’ is added to emphasize that when the flow is non-relativistic, \dot{E} depends explicitly on \dot{M} .

The non-relativistic scalings for the energy loss rate can be compared with those in the relativistic regime. As R_A becomes close to R_L as \dot{M} decreases during the cooling epoch, v_A approaches c and the flow becomes relativistic. At R_L , the ratio of magnetic to kinetic energy density is

$$(5) \quad \Gamma = \left. \frac{B^2}{4\pi\rho c^2} \right|_{R_L} = B_0^2 R_\nu^{2\eta} \Omega^{2\eta-2} c^{1-2\eta} \dot{M}^{-1}$$

and the energy loss rate is

$$(6) \quad \dot{E}_R = -\Gamma \dot{M} c^2 = -B_0^2 R_\nu^{2\eta} \Omega^{2\eta-2} c^{3-2\eta}.$$

For $\eta = 3$ the classical “vacuum dipole” limit is obtained; $\dot{E}_R(\eta = 3) = B_0^2 R_\nu^6 \Omega^4 c^{-3}$ ⁽¹⁾. For the monopole case with $\eta = 2$, \dot{E}_R is larger than the dipole limit by a factor of $c^2/(\Omega R_\nu)^2$ — a factor of ~ 23 for a 10 km MSM with a 1 ms spin period.

The non-relativistic spindown rate is larger than the relativistic spindown rate as a result of mass-loading. To see this explicitly, note that

$$(7) \quad \dot{E}_{\text{NR}}^{\eta=2} / \dot{E}_R^{\eta=2} \simeq 1 B_{0.15}^{-2/3} R_{\nu_{10}}^{-4/3} \dot{M}_{-3}^{1/3} P_1^{2/3},$$

$$(8) \quad \dot{E}_{\text{NR}}^{\eta=3} / \dot{E}_R^{\eta=3} \simeq 8 B_{0.15}^{-6/5} R_{\nu_{10}}^{-18/5} \dot{M}_{-3}^{3/5} P_1^{12/5}.$$

For $\eta = 2$ we see that the ratio is approximately unity, reflecting the fact that for the parameters chosen $R_A \sim R_L$. For slower spin periods, the ratio increases, but

⁽¹⁾ In the true “vacuum dipole” limit \dot{E}_R has a term $\sin \alpha/6$, where α is the angle between the spin axis and the magnetic dipole axis.

not dramatically. In contrast, the ratio of \dot{E}_{NR} to \dot{E}_{R} for the dipole case ($\eta = 3$) is large for MSMs and it has a strong dependence on P . Thus, for a magnetar born with a 10 ms spin period a naive application of the “vacuum dipole” formula underestimates the magnitude of the rotational energy loss rate by a factor of ~ 2000 . More detailed calculations reveal that when $P \sim 10$ ms, for modest neutrino luminosities, \dot{M} is probably closer to $\sim 10^{-5} M_{\odot} \text{ s}^{-1}$ so that the ratio of \dot{E}_{NR} to \dot{E}_{R} is closer to ~ 120 than 2000. Even so, $\dot{E}_{\text{NR}}/\dot{E}_{\text{R}}$ is very large and an application of the relativistic formula in an epoch when \dot{M} is large is incorrect. Of course, if the field structure is not purely dipolar (that is, $\eta < 3$), then the discrepancy between the “vacuum dipole” spindown approximation and the “true” spindown rate becomes even larger. In this case we would compare $\dot{E}_{\text{NR}}^{\eta=2}/\dot{E}_{\text{R}}^{\eta=3} \simeq 25 B_{0.15}^{-2/3} R_{\nu 10}^{-10/3} \dot{M}_{-3}^{1/3} P_1^{8/3}$, a factor of 25 for a MSM. For a magnetar with a 10 ms spin period and lower mass loss rate the ratio becomes $\dot{E}_{\text{NR}}^{\eta=2}/\dot{E}_{\text{R}}^{\eta=3} \simeq 2500 B_{0.15}^{-2/3} R_{\nu 10}^{-10/3} \dot{M}_{-5}^{1/3} P_{10}^{8/3}$. These arguments serve to underscore the fact that in the very early stages of proto-magnetar cooling, the multi-dimensional structure of the wind/magnetic field interaction must be solved consistently (to determine the effective η) and that inferences about the “initial” spin period of magnetars should not be based on an application of the vacuum dipole approximation when \dot{M} is large.

For fixed B_0 , R_{ν} , and Ω the transition from non-relativistic (eq. (4)) to relativistic (eq. (6)) outflow occurs when $R_A \sim R_L$ and this point in time corresponds to a critical mass loss rate $\dot{M}_{\text{crit}} = B_0^2 R_{\nu}^2 \Omega^{2\eta-2} c^{1-2\eta}$, which scales with P^{-2} and P^{-4} for $\eta = 2$ and $\eta = 3$, respectively. Because \dot{M} scales with the neutrino luminosity and because—to first approximation—the luminosity is a monotonically decreasing function of time, this scaling of \dot{M}_{crit} with P implies that the wind is non-relativistic for a larger fraction of the cooling time for longer rotation periods: for $\eta = 3$ and $P = 10$ ms $\dot{M}_{\text{crit}} \simeq 3 \times 10^{-9} M_{\odot} \text{ s}^{-1}$. Such a low \dot{M} may correspond to a time several tens of seconds after collapse.

3. – Are millisecond proto-magnetars GRB central engines?

The spindown timescale $\Omega/\dot{\Omega}$ in the non-relativistic limit is

$$(9) \quad \tau_{\text{JNR}}^{\eta=2} \simeq 30 \text{ s } M_{1.4} \dot{M}_{-3}^{-1/3} R_{\nu 10}^{-2/3} B_{0.15}^{-4/3} P_1^{-2/3},$$

$$(10) \quad \tau_{\text{JNR}}^{\eta=3} \simeq 96 \text{ s } M_{1.4} \dot{M}_{-3}^{-3/5} R_{\nu 10}^{-2/5} B_{0.15}^{-4/5} P_1^{-2/5}.$$

For a MSM with 2×10^{52} ergs of rotational energy we need only wait a fraction of τ_{J} to extract an amount of energy comparable to the supernova energy, $\sim 10^{51}$ ergs. Because τ_{J} (or $\sim \tau_{\text{J}}/5$) is comparable to τ_{KH} we infer that significant spindown may occur during the cooling epoch. As implied by the discussion of \dot{E}_{NR} above, for larger initial spin periods the spindown timescales *decrease*. In the relativistic limit ($R_A \sim R_L$)

$$(11) \quad \tau_{\text{JR}}^{\eta=2} \simeq 34 \text{ s } M_{1.4} R_{\nu 10}^{-2} B_{0.15}^{-2},$$

$$(12) \quad \tau_{\text{JR}}^{\eta=3} \simeq 760 \text{ s } M_{1.4} R_{\nu 10}^{-4} B_{0.15}^{-2} P_1^2.$$

Although $\tau_{\text{JR}}^{\eta=2}$ does not depend on Ω explicitly, it may have an implicit dependence on Ω if the magnetic field is generated by a dynamo [6, 7].

With these timescales in hand we can (in a rudimentary way) attempt to assess the MSM GRB mechanism. A more detailed assessment must await a consistent multi-d MHD solution. From eqs. (3), (4), (9), and (10) we see that in the non-relativistic limit,

regardless of η , the amount of energy extractable on $\sim 10 - 100$ second timescales is in the range appropriate to GRBs. For $\eta < 3$ these conclusions are stronger. Because $\gtrsim 10^{51}$ ergs can be extracted on a timescale shorter than or comparable to the timescale for the SN shock to traverse the progenitor, we expect that the wind may significantly affect the dynamics of the remnant and the ^{56}Ni yield. In this way it may be possible to generate hyper-energetic or 1998bw-like SNe during MSM birth [14, 18]. The inferred energetics and ^{56}Ni yield of SN2003dh and SN1998bw put strong constraints on any GRB mechanism. In the collapsar model a disk wind is thought to generate the ^{56}Ni required to power the SN lightcurve [19, 20]. In the millisecond magnetar model, the energetic wind shocks the material already processed by the supernova shock, perhaps generating the large inferred ^{56}Ni yields; we are currently investigating the timing of this scenario.

Although it seems possible that MSM winds may generate energetic winds at early times, the flow during this mass-loaded wind phase—at least, on average—is not relativistic. It is possible that a strong latitudinal dependence to the mass loss rate may yield relativistic asymptotic velocities for matter emerging from mid-latitudes even when our estimates would indicate $R_A < R_L$, but such a speculation must be tested against realistic multi-d models. It is also possible that large temporal variations in \dot{M} could cause the wind to alternate rapidly between non-relativistic and relativistic. Strong variations in the mass loading could be caused by shearing of large-scale closed magnetic loops on the surface of the fully convective MSM core [16].

As the flow becomes increasingly relativistic (on average), we see from eq. (12) that if the relativistic dipole limit strictly obtains, then the spindown timescale is long for $B \sim 10^{15}$ G and $P \sim 1$ ms. Although relativistic spindown with $\eta = 3$ will affect the asymptotic remnant dynamics by injecting energy over a long timescale, it will probably not generate a GRB with duration ~ 30 seconds and energy $\sim 10^{51}$ ergs. Based on the scalings derived here, higher magnetic field strength, shorter spin period, or $\eta < 3$ is probably required in order for MSM spindown to power GRBs. If one of these possibilities obtains, then from the estimates in ref. [14] we find that essentially all of the magnetic energy at R_L must be transferred to the wind in order for the flow to obtain high asymptotic Lorentz factor with large enough \dot{E} . This is presumably accomplished by magnetic dissipation [21, 22].

4. – Emergence, geometry, and collimation

If a relativistic outflow with the requisite energy to power a GRB can be generated by a MSM, it must emerge from the massive stellar progenitor. The highly energetic non-relativistic wind, which precedes the relativistic outflow, will likely be collimated by hoop stress and will therefore shape the cavity into which the relativistic wind emerges. Because the non-relativistic wind carries little mass in comparison with the overlying star, the relativistic outflow is not additionally hindered in its escape from the progenitor by the preceding slow wind. Hence, if the relativistic outflow can be collimated, then the dynamics of its emergence from the progenitor should be qualitatively similar to models of collapsar jets escaping Type-Ibc progenitors [23].

One important possible objection to the MSM mechanism for GRBs is that it is difficult to collimate relativistic Poynting-flux dominated outflows [24]. Observational evidence for collimation in GRBs is abundant and so, at face-value, this would seem to be a problem. There are at least three responses to this objection. The first possible response is that the interaction between the emerging and energetic (non-relativistic and then relativistic) wind with the overlying post-supernova-explosion ejecta may act to

collimate the outflow. Future multi-d simulations should address this issue in detail. The second potential response is that, in fact, relativistic Poynting-flux dominated winds *can* be efficiently collimated, as in the work of ref. [25]. A final possible response is this: given the basic geometry of the magnetocentrifugal wind, it seems natural to suppose that the asymptotic radial velocity is largest in the equatorial region. This follows both from the fact that centrifugal acceleration is largest at the equator and that the equatorial current sheet may facilitate significant dissipation of the magnetic energy. Pulsars (*e.g.*, the Crab) provide evidence for high Lorentz factor and energetically-dominant equatorial winds [26, 27]. In analogy, is it possible that the geometry of GRBs is “sheet”-like (equatorial) rather than jet-like? The solid angle subtended by a sheet with opening angle θ is $\sim \theta$, whereas for a jet it is $\sim \theta^2$. This fact has implications for the afterglow—a sheet-break rather than a jet-break [28, 29]—and for the GRB census. Inferences from detailed predictions of the flux evolution may be used to rule out or confirm the possibility of a sheet-like geometry in some bursts. I am constructing such models now.

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